

Modification of black-body radiance at low temperatures and frequencies

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Abstract

In contrast to earlier reports, where the spectrum of the *energy density* of photonic black-body radiation modified by SU(2) effects was discussed, we discuss the low-frequency spectrum of the *radiance* at temperatures ranging between 5 and 20 Kelvin. We conclude that compared to the conventional theory the only observable effect is associated with the spectral gap (total screening) which for $T \geq 4.3$ K scales with temperature T as $\frac{\nu^*(T)}{\text{GHz}} = 42.70 \left(\frac{T}{\text{K}}\right)^{-0.53} + 0.21$. We also discuss how a low-temperature black body cavity under the influence of a sufficiently strong static electric field is forced to emit according to Planck's radiation law (pure U(1) theory) even at low frequencies and how this effect can be utilized to measure SU(2) induced deviations.

1 Introduction

The possibility that the $U(1)$ symmetry underlying the propagation of photons as described by the Standard Model of particle physics is the result of a dynamical gauge symmetry breaking of an $SU(2)$ Yang-Mills theory is theoretically [1] and observationally intriguing [2, 3, 4, 5]. In a thermal version of pure and deconfining $SU(2)$ gauge theory associated with a single temperature T this symmetry breaking is induced by the effective thermal ground state which emerges upon a spatial coarse-graining over interacting calorons of topological charge modulus one [6, 7]. The relevant calorons (of trivial holonomy) are stable, periodic, minimal-action, zero energy-momentum, and thus nonpropagating solutions to the Euclidean Yang-Mills equation which possess unit topological-charge modulus [8]. The emerging thermal ground state provides for the temperature-dependent massiveness of two out of the three propagating gauge modes while a third direction of the $SU(2)$ algebra (the photon) remains massless. The fact that slightly below the critical temperature T_c a tiny photon mass due to the Meissner effect induced by a condensate of (electric) monopoles seems to be implicit in the low-frequency CMB data, see [9] and references therein and [10] for an interpretation, fixes T_c to be very closely above the present CMB baseline temperature of about 2.73 Kelvin.

In [11, 12, 13, 14] an account of the computation of radiative corrections in the effective theory has been given. In particular, the screening function G for thermalized photon propagation in the deconfining phase was computed exactly on the one-loop level¹ in [14]. Our predictions for the modified black-body spectrum that rely on function G so far were for the spectral energy density (usually referred to as ‘spectral intensity’ in [4, 12, 14]) which is not directly accessible experimentally². For bolometry and radiometry the relevant quantity is the spectral radiance L .

The purpose of this note is an investigation of the characteristics of the radiance spectrum at low temperature and frequency as well as a discussion of observable effects induced by $SU(2)$ gauge dynamics.

¹This is more than sufficient for any practical purpose, see [15].

²Academically, it is in principle accessible through gravitational interactions.

2 Bolometry and radiometry of SU(2) photons

Here we derive the photonic radiance of an SU(2) Yang-Mills theory subject to a modified dispersion relation at low temperatures and frequencies [14] in comparison with the conventional U(1) Planck spectrum. In contrast to earlier publications, where natural units were used [4, 12, 14], we exclusively work in SI units in this report. Recall that the spectral energy density $I(\omega)$ of a photon gas in thermal equilibrium is given by the number of modes per volume available within a frequency interval times the average (thermal) energy per mode. One has

$$\begin{aligned} dE &= I(\omega)d\omega = 2 \frac{d^3\mathbf{k}}{(2\pi)^3} \times \frac{\hbar\omega}{e^{\frac{\hbar\omega}{kT}} - 1} \\ &= \frac{4}{h^2} \frac{\omega}{e^{\frac{\hbar\omega}{kT}} - 1} \mathbf{p}^2(\omega) \frac{dp}{d\omega} d\omega \\ \Rightarrow I(\nu) &= \frac{8\pi}{h^2} \frac{\omega}{e^{\frac{\hbar\omega}{kT}} - 1} \mathbf{p}^2(\omega) \frac{dp}{d\omega}, \end{aligned} \tag{1}$$

where k is Boltzmann's constant, h is Planck's quantum of action ($\hbar = \frac{h}{2\pi}$), \mathbf{p} denotes a photon's spatial momentum, $\omega = 2\pi\nu$, and T is the absolute temperature. The factor $\frac{dp}{d\omega}$ in Eq. (1) accounts for the dispersion relation $\omega = \omega(|\mathbf{p}|)$. In a U(1) theory this relation is given by $(\hbar\omega)^2 = (cp_0)^2 = c^2\mathbf{p}^2$ with c the speed of light in vacuum. For the photons of an SU(2) Yang-Mills theory the dispersion law is substantially altered at low frequencies and low temperatures, an effect which is described by the (anti)screening function³ G . Namely, one has

$$p_0^2 - \mathbf{p}^2 = G(|\mathbf{p}|, T). \tag{2}$$

For photons to propagate it must hold that $p_0^2 \geq G$ (otherwise $|\mathbf{p}|$ becomes imaginary). This restricts the possible values of ω associated with energy transport to $\omega \geq \omega^*$, with ω^* defined as the root of $|\mathbf{p}(\omega^*)| = 0$.

In conventional U(1) theory both quantities, spectral energy density $I(\nu)$ and spectral radiance $L(\nu)$, are proportional to one another, $L(\nu) = \frac{c}{4\pi} \times I(\nu)$ [16]. In a deconfining SU(2) plasma, however, c must be replaced by the

³Microscopically, this screening or antiscreening takes place by monopole-antimonopole creation or successive scattering of photons off monopoles and antimonopoles, respectively.

photon's group velocity v_g defined as

$$v_g \equiv \partial_{|\mathbf{p}|} E = \hbar \partial_{|\mathbf{p}|} \omega = \partial_{|\mathbf{k}|} \omega, \quad (3)$$

where $E = \hbar \nu = c p_0$ is photonic energy and $\mathbf{k} \equiv \mathbf{p}/\hbar$ the wave number vector. Thus in calculating the spectral radiance $L(\nu)$ for SU(2) photons, the factor $\frac{dp}{d\omega}$ cancels the one in Eq. (1), and we obtain

$$\begin{aligned} L_{\nu}^{\text{SU}(2)}(T, \nu) &= \frac{2h}{c^2} \frac{\nu^3}{e^{\frac{\hbar\nu}{kT}} - 1} \times \left(1 - \frac{c^2 G}{(h\nu)^2}\right) \theta(\nu - \nu^*) \\ &= L_{\nu}^{\text{U}(1)} \times \left(1 - \frac{c^2 G}{(h\nu)^2}\right) \theta(\nu - \nu^*) \end{aligned} \quad (4)$$

where $L_{\nu}^{\text{U}(1)}$ denotes the Planckian spectral radiance.

As shown in [14], the (anti)screening function G may be calculated in a selfconsistent way using numerical methods. To make contact with the real world the critical temperature T_c for the deconfining-preconfining phase transition, which is the only free parameter of a thermalized SU(2) quantum Yang-Mills theory, must be determined experimentally. In [10] we have given observational reasons why T_c should be very closely above the baseline temperature of about 2.73 Kelvin of the present cosmic microwave background.

After fixing T_c the modified radiance spectra for SU(2) photons may be calculated for various physical temperatures. In Figs. 1, 2, 3 we show results for $T = 5.4, 8.0$, and 11.0 Kelvin, respectively. The red lines correspond to the calculated SU(2) radiance, the conventional U(1) Planck spectrum is depicted in grey. Notice the regime of total screening (suppression of spectral radiance down to zero) and the cross-over to a regime of slight antiscreening (excess of spectral radiance) in all cases. Fig. 4 shows the difference in spectral radiance $\Delta L_{\nu}(T, \nu)$, defined as

$$\Delta L_{\nu}(T, \nu) = L_{\nu}^{\text{SU}(2)}(T, \nu) - L_{\nu}^{\text{U}(1)}(T, \nu) \quad (5)$$

for $T = 5.4$ K.

A type of bolometric experiment can be conceived as follows. Let the apertures of an isolated low-temperature U(1) black body at temperature T_1 and that of an SU(2) black body of identical characteristics at temperature T_2 face each other, and exchange radiant energy. Have the SU(2) black body be linked to a large heat reservoir to keep its wall temperature T_2 constant. Furthermore, switch in a band-width filter within the common aperture tuned

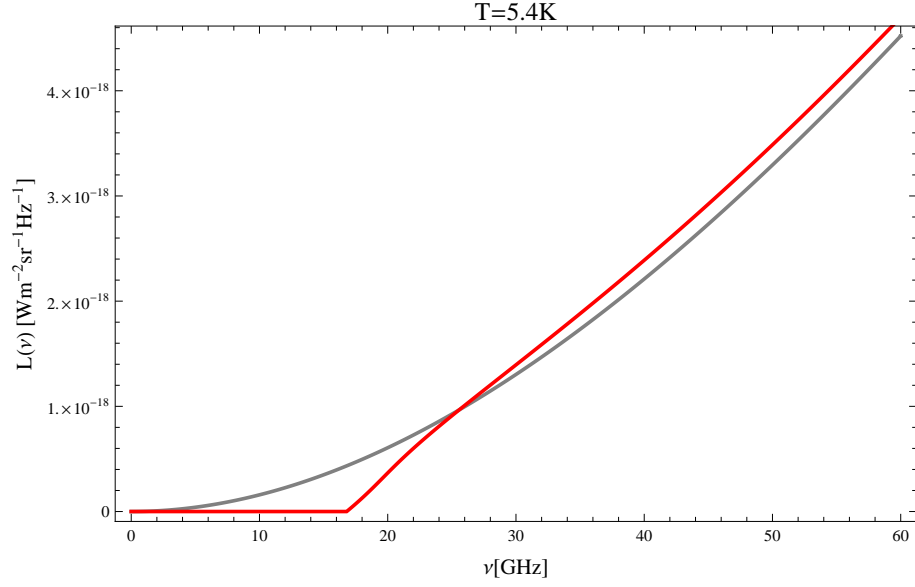


Figure 1: Comparison between the black body spectral radiances for $T = 5.4$ K of an SU(2) (red) and a U(1) (grey) theory.

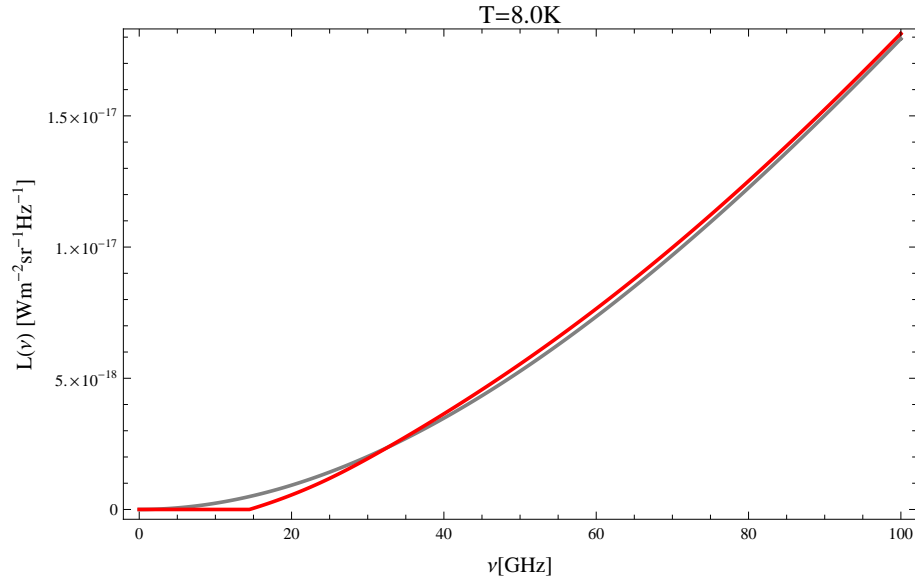


Figure 2: Comparison between black body spectral radiances for $T = 8.0$ K of an SU(2) (red) and a U(1) (grey) theory.

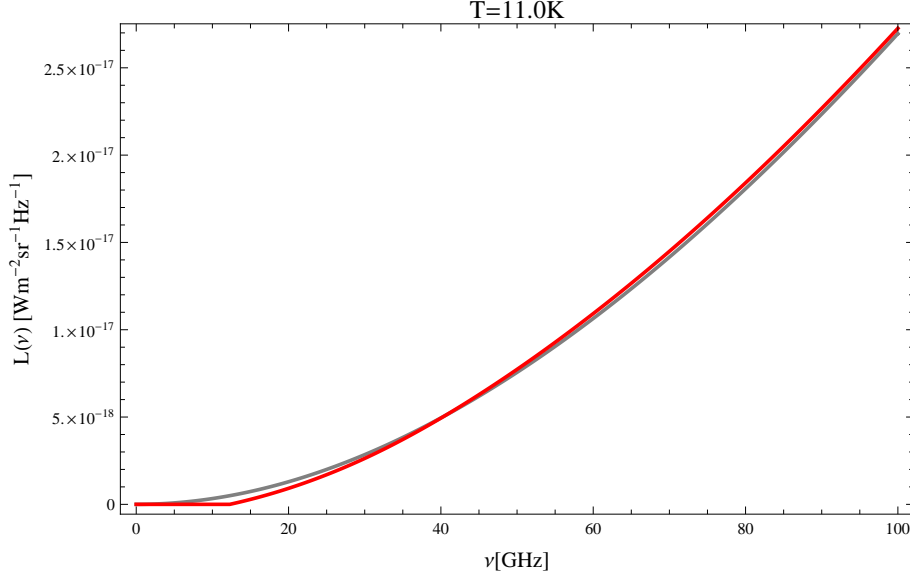


Figure 3: Comparison between black body spectral radiances for $T = 11.0$ K of an SU(2) (red) and a U(1) (grey) theory.

to the region of the SU(2) black-body gap: Such a filter absorbs photons above frequency ν^* no matter which cavity they come from⁴. U(1) photons within the SU(2) spectral gap are absorbed by the SU(2) thermal ground state [17]: They create unresolvable monopole-antimonopole pairs. Thus a small amount of energy per time flows across the common aperture from the U(1) cavity to the SU(2) cavity, and zero radiation temperature within the U(1) cavity will asymptotically be reached (radiation refrigeration).

For $T_1 = T_2$ and blocking off the gap region $0 \leq \nu \leq \nu^*(T_1)$ by a complementary band-width filter, the regime of propagating photon frequencies would not allow for any heat exchange since $L^{\text{SU}(2)}(\nu, T_2) d\Omega_{\text{SU}(2)}$ and $L^{\text{U}(1)}(\nu, T_1 = T_2) d\Omega_{\text{U}(1)}$ do match. Here $d\Omega_{\text{SU}(2)}$ and $d\Omega_{\text{U}(1)}$ are solid-angle elements defined on a sphere centered at a point within the common aperture.

⁴ From the three effective SU(2) gauge modes only the tree-level massless mode (the photon) interacts with electric charges and so can be absorbed by the material in the band-width filter. If a propagating photon emerges from the SU(2) cavity then the absorbed power per frequency interval is identical to that of a propagating photon stemming from the U(1) cavity because the additional factor $\frac{1}{1 - \frac{c^2 G}{(h\nu)^2}}$ in the SU(2) spectral radiance, see Eq. (4), is canceled, see discussion below Eq. (13).

By Snell's law [18] we have

$$\frac{d\Omega_2}{d\Omega_1} = \frac{v_{ph}^2}{c^2}, \quad (6)$$

where the phase velocity v_{ph} of energy propagation inside the SU(2) black-body is given by $v_{ph} \equiv \frac{\omega}{k} = \frac{c}{\sqrt{1 - \frac{c^2 G}{(h\nu)^2}}}$, and therefore

$$L^{\text{SU}(2)}(\nu, T_2) \frac{d\Omega_{\text{SU}(2)}}{d\Omega_{\text{U}(1)}} = L^{\text{U}(1)}(\nu, T_1 = T_2). \quad (7)$$

A potentially interesting quantity for experiments is the *radiance U(1) line temperature* $T_P(\nu)$. The quantity $T_P(\nu)$ is the temperature a conventional U(1) black body must possess in order to reach the following (bolometric) equilibrium condition [18], compare with Eq. (6):

$$L_\nu^{\text{U}(1)} = L_\nu^{\text{SU}(2)} \times \frac{1}{1 - \frac{c^2 G}{(h\nu)^2}}. \quad (8)$$

One has

$$T_P(L_\nu^{\text{SU}(2)}(\nu, T)) \equiv \frac{h\nu}{k} \frac{1}{\ln \left[\frac{2h}{c^2} \left(1 - \frac{c^2 G}{(h\nu)^2} \right) \frac{\nu^3}{L_\nu^{\text{SU}(2)}(\nu, T)} + 1 \right]}. \quad (9)$$

Again, on the right-hand side of Eq. (8) the factor $\left(1 - \frac{c^2 G}{(h\nu)^2} \right)$, which apart from the factor $\theta(\nu - \nu^*)$ distinguishes SU(2) from U(1) spectral radiance, see Eq. (4), is canceled. Therefore T_P is the same as the SU(2) wall temperature T_2 except for the regime of total screening where $T_P \equiv 0$. That is, no heat is effectively exchanged by frequencies above the spectral gap according to the bolometric equilibrium condition (8), and within the spectral gap no U(1) photons are required in the balance of Eq. (9). An essential question, addressed in Sec. 3, is how the above-assumed hypothetic low-temperature U(1) black body can be realized experimentally.

For completeness let us give some characterization of the factor

$$\left(1 - \frac{c^2 G}{(h\nu)^2} \right) \theta(\nu - \nu^*)$$

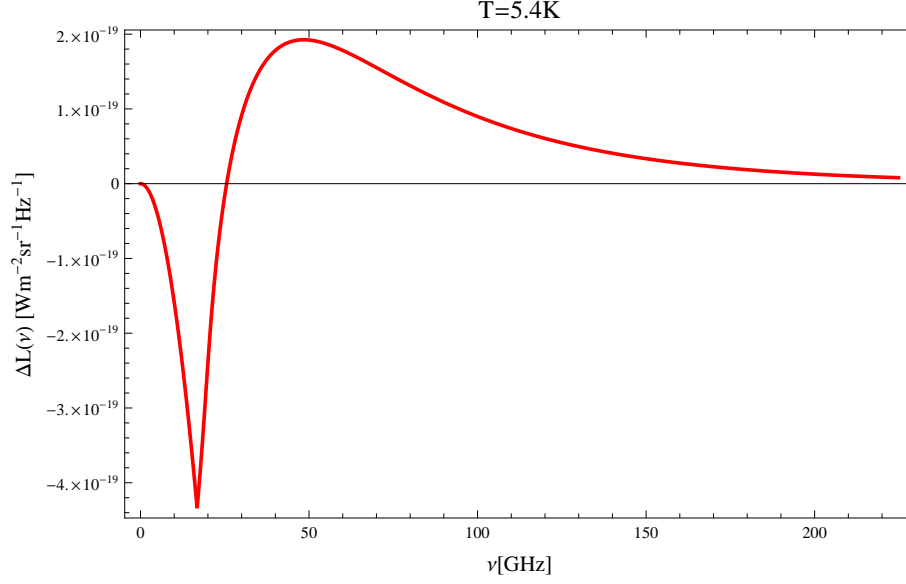


Figure 4: The difference in spectral radiance $\Delta L(\nu, T)$ at $T = 5.4$ K as a function of frequency ν .

which converts U(1) to SU(2) spectral radiance and comprises of the characteristic frequencies ν^* , ν_c , ν_M which are implicitly defined as follows:

$$\begin{aligned}
 |\mathbf{p}|(\nu^*) &= 0, \\
 G(\nu_c, T) &= 0, \quad \text{and} \\
 \frac{G(\nu_M, T)}{\nu_M^2} &= \min\left\{\frac{G(\nu, T)}{\nu^2}\right\}.
 \end{aligned} \tag{10}$$

Lowering ν at fixed T , the points ν^* , ν_c , and ν_M describe the onset of total screening (no photon propagation), the cross-over between screening and antiscreening ($G = 0$), and the maximal antiscreening ($G < 0$), respectively. For $T > 8$ K the critical points ν_c , ν_M and ν^* were numerically fitted to a

power law in T . We obtain the following results:

$$\begin{aligned}\frac{\nu_c(T)}{\text{GHz}} &= 1.83 \left(\frac{T}{\text{K}}\right)^{1.12} + 13.48 \\ \frac{\nu_M(T)}{\text{GHz}} &= 3.45 \left(\frac{T}{\text{K}}\right)^{1.08} + 12.90 \\ \frac{\nu^*(T)}{\text{GHz}} &= 42.70 \left(\frac{T}{\text{K}}\right)^{-0.53} + 0.21.\end{aligned}\tag{11}$$

Fig. 5 shows calculated points overlaid with the fitted curves.

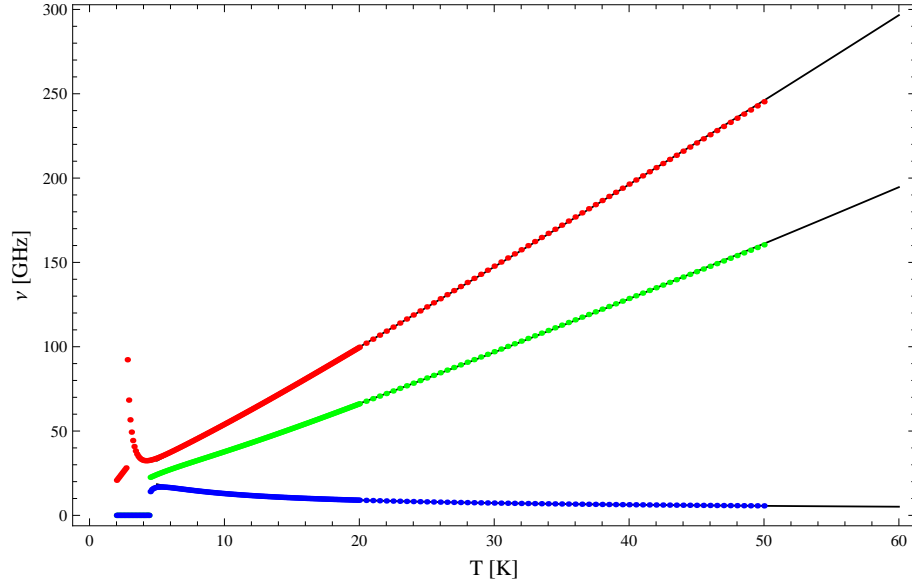


Figure 5: Plots of the T dependence of the characteristic points of $L_\nu^{SU(2)}(T, \nu)$. The blue points depict $\nu^*(T)$, the green points $\nu_c(T)$, and the red points $\nu_M(T)$.

We now consider a radiometric approach by placing an antenna inside the $SU(2)$ plasma. Independent of (lossless) propagation properties, one arrives at the following expression for the power $P_{\Delta\nu}(\nu)$ within band width $\Delta\nu$ absorbed by the antenna [16]:

$$P_{\Delta\nu}(\nu) = \theta(\nu - \nu^*) \int_{\nu}^{\nu + \Delta\nu} d\nu' \frac{h\nu'}{\exp\left[\frac{h\nu'}{kT}\right] - 1}.\tag{12}$$

For $h\nu \ll kT$, which is certainly the case for the spectral range we are interested in, Eq. (12) simplifies as

$$P_{\Delta\nu}(\nu) \approx \theta(\nu - \nu^*) \Delta\nu kT. \quad (13)$$

Apart from the θ -function prefactor in Eqs. (12) and (13) the expressions are identical to the U(1) situation since a factor $|\mathbf{k}|^2$ in $L_\nu^{SU(2)}(\nu, T)$ is cancelled by a factor $|\mathbf{k}|^{-2}$, see [16]. Radiometric measurements inside the predicted screening regime may therefore be able to test for the presence of a fundamental SU(2) ground state resulting from gauge dynamics.

3 Preparation of a U(1) black body at low temperatures

The concept of a U(1) line temperature relies on one's ability to prepare a black-body source in such a way that it emits according to Planck's radiation law even at small temperatures where we expect modifications thereof. Since a homogeneous electric field of energy density well above that of the thermal ground state effectively switches off the indirect coupling of the photon to the ground state via scattering processes involving the massive vector modes a black-body cavity with the bulk of its volume transcended by such an electric field effectively acts as a U(1) emitter. In the single-photon counting experiment of Ref. [19] the (two-step laser) excitation of the $111_{s_{1/2}}$ Rydberg state of ^{85}Rb atoms towards the $111_{p_{3/2}}$ state by absorption of thermal photons of frequency 2527 MHz, prepared within a tunable cavity, can be exploited to learn about the temperature dependence of the mean photon number $\bar{n}(T)$ at this frequency. The measurement was performed for temperatures T ranging from 67 mK up to 1 K. As discussed in [19], no deviation of $\bar{n}(T)$ from the U(1) expected Bose-Einstein distribution was observed. It is important to note that a static, electric stray field of $|\vec{E}| \sim 25 \text{ mV/cm}$ was present in the cavity during the experiment.

Such an electric field, however, would cause a sizable distortion of the thermal SU(2) ground state which is sufficient to render the system effectively U(1). The (unresolvable) electric monopoles residing in the thermal ground state are accelerated by the external field in a parallel or antiparallel way, thus acquire kinetic energy which they subsequently disperse by collisions thereby increasing the energy density of the thermal ground state to an

effective temperature largely disparate to the temperature of the radiation which is kept in thermal equilibrium with the cavity walls. This separation of excitation from ground state physics frees the propagation of photons from any ground-state induced effect and thus renders their dispersion law trivial, that is, of the conventional U(1) type. This effect may be used to test whether thermalized but otherwise unadulterated photon propagation that we observe in Nature really is a manifestation of strongly interacting thermal SU(2) gauge dynamics.

The energy density ρ_E of the external electric field is given as

$$\rho_E = \frac{\epsilon_0}{2} \vec{E}^2, \quad (14)$$

where in SI units $\epsilon_0 = 8.8542 \cdot 10^{-12} \text{ J}/(\text{Vm}^2)$. The energy density of the SU(2) thermal ground state ρ_{gs} in SI units is given as

$$\rho_{gs} = 4\pi \Lambda_{CMB}^3 \frac{k_B}{(\hbar c)^3} T, \quad (15)$$

where $\Lambda_{CMB} = 2\pi \frac{k_B T_c}{\lambda_c}$, $T_c = 2.725 \text{ K}$ [10], $\lambda_c = 13.87$ [1], $k_B = 1.3807 \cdot 10^{-23} \text{ J/K}$, $c = 2.9979 \cdot 10^8 \text{ m/s}$, $\hbar = \frac{6.6261}{2\pi} \cdot 10^{-34} \text{ Js}$. Setting $\rho_E = \rho_{gs}$ at $|\vec{E}| \sim 25 \text{ mV/cm}$ [19], we obtain an effective ground-state temperature of $\sim 10^3 \text{ K}$. Thus the energy density of the external electric field is more than three orders of magnitude larger than that of the SU(2) thermal ground at 1 K and below. Thus we may safely assume a decoupling of the ground-state physics from the propagation properties of photons. In a rough estimate on an admissible field strength to not distort the SU(2) ground-state physics sizably ρ_E should be less than ρ_{gs} . For example, demanding that $\rho_E \sim 0.1 \rho_{gs}$ at $T = 5.4 \text{ K}$ implies an electric field strength of about $|\vec{E}| \sim 0.2 \text{ mV/cm}$. Thus to produce a U(1) black body at low temperature one should work with a much larger value of $|\vec{E}|$. A more refined treatment of the effects induced on the thermal ground state by external fields, considered as small perturbations to undistorted SU(2) Yang-Mills thermodynamics, should be performed within the realm of linear-response theory. We leave this to future investigation.

4 Summary and Conclusions

In this note we have investigated the experimental consequences of the assumption that, fundamentally, photon propagation is described by an SU(2)

rather than U(1) gauge principle. We have considered bolometric and radiometric methods, and we have shown that the only region where a differences to the conventional theory can be predicted is associated with the spectral gap $0 \leq \nu \leq \nu^*(T)$ with the T dependence of ν^* given in Eq.(11) (total screening). We also have elucidated how a U(1) black body can be prepared at low temperature by virtue of decoupling its thermal ground state from its excitations due to the application of a static, homogeneous electric field.

Acknowledgments

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